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## Supersymmetry Breaking by Hidden Matter Condensation in Superstrings

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### ABSTRACT

We show that supersymmetry can be broken mainly by hidden matter condensates in the observable matter direction in generic superstring models. This happens only when the fields whose VEVs give masses to hidden matter do not decouple at the condensation scale. We find how the parameters of the string model and the vacuum determine whether supersymmetry is broken mainly by hidden matter or gaugino condensates and in the matter or moduli directions.

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## 1. Introduction

One of the most important but least understood aspects of superstring theories is supersymmetry (SUSY) breaking. It is well-known that SUSY must be broken non-perturbatively and around the  $TeV$  scale in the observable sector (to solve the hierarchy problem). Due to our lack of understanding of non-perturbative string effects, the best we can do is to investigate SUSY breaking by non-perturbative phenomena in field theories which are low-energy (i.e.  $E \ll M_P$ ) limits of superstring theories (such as string induced supergravity [1]).

The most common way of achieving dynamical SUSY breaking in superstrings is by hidden gaugino condensation in supergravity [2] theories which are obtained from the massless sector of superstrings. In this scenario, hidden sector gaugino condensates form when a non-Abelian hidden gauge group becomes strong at a hierarchically small scale,  $\Lambda_H \ll M_P$ . The presence of such hidden sectors with non-Abelian gauge groups is a generic feature of superstring models. The condensation is taken into account by a non-perturbative superpotential,  $W_{np}$ , [3] which has all the required symmetry properties. One then finds that SUSY can be broken in the moduli direction and in a phenomenologically acceptable way (with the well-known problems of the vanishing cosmological constant and the stability of the dilaton potential).

Most hidden sectors of superstring models also contain hidden matter in the vector representations of the hidden gauge groups which condense with the gauginos. Their presence not only affects the running of the hidden gauge coupling constant but also modifies  $W_{np}$  [3,4]. In addition, matter condensation can also be the source of SUSY breaking. Surprisingly, SUSY breaking by hidden matter condensation has not attracted much attention until recently [5]. In Ref. (5), the effect of hidden matter condensation on  $F$  terms was examined in the framework of a realistic string model by examining the hidden matter mass terms. In the following we will often refer to the model of Ref. (5) as a concrete example of our results. In this letter, we generalize the results of Ref. (5) by including the effects

of gaugino condensation and the full-fledged  $W_{np}$  for superstrings. We argue that SUSY can be broken by hidden matter rather than hidden gaugino condensates and in the observable matter direction rather than the moduli directions. We show that this is a realistic possibility under quite generic conditions if the fields whose VEVs give masses to hidden matter do not decouple at the condensation scale. Our aim here is only to show the possibility of this new kind of SUSY breaking in generic superstring models. Whether this is the case or not in a specific string model depends on the details of the model such as the hidden sector gauge group and matter, the hidden matter mass terms etc. as we will show below. This can only be investigated in the framework of a specific model and with a detailed numerical analysis of the scalar potential which we defer to the future.

## 2. SUSY breaking by hidden gaugino condensation

In this section we briefly review the gaugino condensation scenario in superstrings. We consider a superstring model with a number of generic properties to be outlined below rather than a specific one (such as standard-like superstring models [6]) for two reasons. First this makes the discussion about SUSY breaking in realistic superstring models more general. Second the properties we outline below (and in the next section) can be seen loosely as necessary conditions for supersymmetry breaking by hidden matter condensation. We consider a superstring model in the four dimensional free fermionic formulation [7] with the following properties:

(a) The massless spectrum of the superstring model is divided into observable and hidden sectors. The observable sector contains a large number of states ( $\phi_i$ ) which are Standard Model (SM) singlets coming from the Neveu–Schwarz and some twisted sectors in addition to the chiral generations.

(b) The hidden sector contains one (or more)  $SU(N)$  non–Abelian gauge group(s) with  $M$  copies of matter ( $h_i, \bar{h}_i$ ) in the vector representations  $N + \bar{N}$ . The hidden matter states obtain masses from non–renormalizable terms,  $W_n$ , in the superpotential. Thus, the hidden matter mass matrix is non–singular and the

SUSY vacuum is stable [8].  $M < 3N$  so that the hidden gauge group is asymptotically free and condenses at the scale  $\Lambda_H \sim M_v \exp(8\pi^2/bg^2)$  where  $b = M - 3N$ .

(c) The Kahler potential is generically given by [9]

$$K(S, S^\dagger, T, T^\dagger, \phi_i, \phi_i^\dagger) = -\log(S + S^\dagger) - 3\log(T + T^\dagger) - \sum_i \phi_i \phi_i^\dagger, \quad (1)$$

where  $S, T$  and  $\phi_i$  are the dilaton, (overall) modulus and matter fields respectively. These fields are in the “supergravity basis” and are related to the massless string states by well-known transformations. (For a recent discussion of moduli and Kahler potentials in free fermionic models see Ref. (10).)

(d) The string vacuum is supersymmetric at the Planck scale,  $M_P$ . This is guaranteed by satisfying the F and D constraints obtained from the cubic superpotential  $W_3$  (which is trilinear in  $\phi_i$  and  $h_i$ ) and the local charges of the states. As we will see below,  $W_3$  does not get any higher order corrections as long as the hidden gauge group does not condense at  $\Lambda_H \ll M_P$ . Therefore,  $W_3$  is the exact superpotential until hidden sector condensation which results in SUSY breaking. The set of F and D constraints is given by the following equations [11]:

$$D_A = \sum_i Q_i^A |\langle \phi_i \rangle|^2 = \frac{-g^2 e^{\phi_D}}{192\pi^2} \text{Tr}(Q_A) \frac{1}{2\alpha'}, \quad (2a)$$

$$D^j = \sum_i Q_i^j |\langle \phi_i \rangle|^2 = 0, \quad (2b)$$

$$\langle W_3 \rangle = \left\langle \frac{\partial W_3}{\partial \phi_i} \right\rangle = 0, \quad (2c)$$

where  $\phi_i$  are the matter fields and  $Q_i^j$  are their local charges.  $\alpha'$  is the string tension given by  $(2\alpha')^{-1} = g^2 M_P^2 / 32\pi = g^2 M_v^2$  and  $\text{Tr}(Q_A) \sim 100$  generically in realistic string models. Eq. (2a) is the D constraint for the anomalous  $U(1)_A$  which is another generic feature of realistic string models [11]. We see that some SM singlet scalars must get Planck scale VEVs of  $O(M_v/10)$  in order to satisfy Eq. (2a) and preserve SUSY around the Planck scale. Then, due to the other F and D constraints most of the other SM singlet scalars also obtain VEVs of  $O(M_v/10)$ .

(e) There are non-renormalizable (order  $n > 3$ ) terms in the superpotential which are generically of the form

$$W_n = c_n g^{n-2} h_i \bar{h}_j \phi_{j_1} \phi_{j_2} \dots \phi_{j_{n-2}} \eta(T)^{2n-6} M_v^{3-n}, \quad (3)$$

which are obtained from the world-sheet correlators  $A_n \sim \langle V_1^f V_2^f V_3^b \dots V_n^b \rangle$  using the rules of Ref. (12).  $c_n$  are numerical coefficients of  $O(1)$  and  $\eta$  is the Dedekind eta function. The powers of  $\eta$  and  $M_v$  are such that the term  $W_n$  has modular weight  $-3$  and dimension 3 as it should. ( $\phi_i$  and  $h_i$  have modular weight  $-1$  which is a generic feature of realistic models in the free fermionic formulation.) These terms contain both observable and hidden sector states. Once the fields  $\phi_i$  get VEVs, they give masses to the hidden states  $h_i, \bar{h}_i$ . Therefore all the  $n > 3$  terms in Eq. (3) can be seen as hidden matter mass terms. These corrections to  $W_3$  when they become non-zero, (i.e. when hidden matter condensates  $\Pi_{ij} = h_i \bar{h}_j$  form) give corrections to the cubic level F constraints in Eq. (2a) and destabilize the original SUSY vacuum as was shown in Ref. (5). (In general, there can also be terms of the form  $c_n \phi_{i_1} \phi_{i_2} \dots \phi_{i_n}$ , i.e. non-renormalizable terms with only observable fields. These vanish in standard-like models [5] and we assume that they are not present in the following.)

The above assumptions about the string model are relatively mild since they are all generic features of realistic superstring models such as standard-like models [6]. We investigate what happens when the hidden gauge group condenses at  $\Lambda_H$  in two different cases: (a) when  $\phi_i$  are heavy, i.e.  $m_{\phi_i} \gg \Lambda_H$  and they decouple at  $\Lambda_H$ , and (b) when  $\phi_i$  are light i.e.  $m_{\phi_i} < \Lambda_H$  and they remain in the spectrum. In both cases we assume that the hidden matter states  $h_i, \bar{h}_i$  do not decouple from the spectrum at  $\Lambda_H$  (otherwise obviously there can only be gaugino condensation).

Case (a) is the case previously investigated in superstring models [4]. When the hidden gauge group condenses at  $\Lambda_H$ , gaugino condensates  $Y^3$  and matter condensates  $\Pi_{ij} = h_i \bar{h}_j$  form. The non-perturbative effective superpotential obtained

from the Ward identities and modular invariance is

$$W_{np} = \frac{1}{32\pi^2} Y^3 \log\{ \exp(32\pi^2 S) [c\eta(T)]^{6N-2M} Y^{3N-3M} \det \Pi \} - \text{tr} A \Pi, \quad (4)$$

where  $c$  is a constant and  $A$  is the hidden matter mass matrix given by the  $n > 3$  terms in Eq. (3). The last term corresponds to the sum of all the  $n > 3$  terms in Eq. (3). The observable matter fields  $\phi_i$  appear only in the mass matrix  $A$ . In the flat limit  $M_P \rightarrow \infty$ , gravity decouples and one gets a globally SUSY vacuum at which (in addition to Eqs. (2a-c))

$$\frac{\partial W_{tot}}{\partial Y} = \frac{\partial W_{tot}}{\partial \Pi} = 0, \quad (5)$$

where  $W_{tot} = W_3 + W_{np}$ . We can replace  $W_{tot}$  in Eq. (5) by  $W_{np}$  since  $W_3$  does not contain  $Y^3$  or  $\Pi$ . The  $n > 3$  terms,  $W_n$  which are the hidden matter mass terms, are already included in  $W_{np}$  through  $\text{tr} A \Pi$ . The solutions to Eq. (5) are used to obtain the composite fields  $Y^3$  and  $\Pi$  in terms of  $S, T, A$

$$\frac{1}{32\pi^2} Y^3 = (32\pi^2 e)^{M/N-1} [c\eta(T)]^{2M/N-6} [\det A]^{1/N} \exp(-32\pi^2 S/N), \quad (6)$$

and

$$\Pi_{ij} = \frac{1}{32\pi^2} Y^3 A_{ij}^{-1}. \quad (7)$$

Eqs. (6) and (7) are used to eliminate the composite fields in  $W_{np}$  and then

$$W_{np}(S, T) = \Omega(S) h(T) [\det A]^{1/N}, \quad (8)$$

where

$$\Omega(S) = -N \exp(-32\pi^2 S/N), \quad (9a)$$

$$h(T) = (32\pi^2 e)^{M/N-1} [c\eta(T)]^{2M/N-6}. \quad (9b)$$

In  $W_{np}$  all the information about the matter condensates,  $\Pi$ , and the observable fields  $\phi_i$  is contained in the term  $\det A$ . When  $m_{\phi_i} \gg \Lambda_H$  and  $\phi_i$  decouple,

one simply substitutes the VEVs  $\langle\phi_i\rangle$  obtained from the solution to the F and D constraints in  $\det A$ .  $\phi_i$  are longer dynamical fields since at the scale  $\Lambda_H$  these heavy fields cannot be excited but simply sit at their VEVs. In this sense,  $\phi_i$  are similar to the composite fields  $Y^3$  and  $\Pi$  which are also eliminated from  $W_{np}$ . All  $\phi_i$  do is to give masses to the hidden matter states  $h_i, \bar{h}_i$  through their VEVs. As a result, in this case the only effect of matter condensates  $\Pi_{ij}$  is to change the scale of the gaugino condensate  $Y^3$  through  $\det A$ .

It is well-known that  $W_{np}$  above breaks SUSY in the modulus direction (but not in the dilaton direction), i.e.  $\langle F_T \rangle \neq 0$  (but  $\langle F_S \rangle = 0$ ) where [1]

$$F_k = e^{K/2}(W_k + K_k W). \quad (10)$$

The subscript denotes differentiation with respect to fields and  $k = S, T$ . Here  $W = W_{tot}$ . The vacuum is obtained by minimizing the scalar potential [1]

$$V = \sum_k |F_k|^2 (G_k^k)^{-1} - 3e^K |W|^2, \quad (11)$$

where  $k = S, T$  and  $G = K + \log|W|^2$ . We do not give explicit expressions for  $F_k$  since they are special cases of the ones we obtain in the next section where we include the effects of matter condensation and observable fields  $\phi_i$  with  $m_{\phi_i} < \Lambda_H$ .

### 3. SUSY breaking by hidden matter condensation

In this section we consider case (b) mentioned above in which  $m_{\phi_i} < \Lambda_H$  and  $\phi_i$  remain in the spectrum. Then,  $\phi_i$  should be treated as dynamical fields similar to  $S$  and  $T$  since they can be excited due to their small masses. Now  $W = W(S, T, \phi_i)$  where from Eq. (8) all the  $\phi_i$  dependence is in the term  $\det A$  which arises due to the matter condensates  $\Pi_{ij}$ . As a result, in addition to  $F_{S,T}$  one should also check whether  $F_{\phi_i}$  vanishes or not in the vacuum. Also, it may now be possible to break SUSY mainly by hidden matter condensation rather than hidden gaugino condensation.

$\det A$  is a product of mass terms given generically by Eq. (3). Thus without any loss of generality, we can assume that it has the form

$$\det A = k S^{-r} \phi_i^{s_i} \eta(T)^t \quad r, s, t > 0, \quad (12)$$

where the  $S$  dependence is obtained from the relation  $g^2 = 1/S$  (at the string tree level and for level one Kac–Moody algebras).  $\phi_i$  denotes any matter field which appears in  $\det A$  and  $s_i$  is its power.  $k$  is a constant of  $O(1)$  which is given by the product of the relevant  $c_n$  in Eq. (3). In fact, this is the form of  $\det A$  which was obtained from the explicit model of Ref. (5) with  $r = 7$ ,  $t = 22$  and  $s_i = 1, 5$  depending on the field  $\phi_i$ . (In general,  $\det A$  is a sum of terms like that in Eq. (12).) We see that there is a new  $S$  and  $T$  dependence in  $W_{np}$  due to  $\det A$ . Taking this into account, we find for the  $F$  term in the dilaton direction

$$F_S = \frac{e^{-\phi_i \phi_i^\dagger / 2}}{(S + S^\dagger)^{1/2} (T + T^\dagger)^{3/2}} h(T) [\det A]^{1/N} \times \left\{ \Omega_S - \frac{\Omega}{(S + S^\dagger)} + \Omega(\log[\det A]^{1/N})_S \right\}. \quad (13)$$

The first two terms in the curly brackets are the usual ones coming from gaugino condensation. The last term gives the contribution of the matter condensates (through  $\det A$ ) to  $F_S$ . Assuming the above form for  $\det A$  we get

$$\frac{\partial(\log[\det A]^{1/N})}{\partial S} = \frac{-r}{NS}. \quad (14)$$

Using Eq. (9a) for  $\Omega(S)$  and the fact that  $S \sim 1/2$  in order to have gauge coupling unification around  $10^{18}$  GeV, we find that the first term in the curly brackets always dominates the other two for realistic values of  $r$  and  $N$ . For example, in the explicit example of Ref. (5),  $N = 5$  and  $r = 7$  and therefore the gaugino part is larger than the matter part by a factor of  $\sim 100$ . In other words, the effect of matter condensates on  $F_S$  is negligible.



For the  $F$  term in the modulus direction we find

$$F_T = \frac{e^{-\phi_i \phi_i^\dagger/2}}{(S + S^\dagger)^{1/2}(T + T^\dagger)^{3/2}} \Omega(S) [\det A]^{1/N} \times \left\{ h_T - \frac{3h}{(T + T^\dagger)} + h(\log[\det A]^{1/N})_T \right\}. \quad (15)$$

As for  $F_S$ , the first two terms in the curly brackets arise from gaugino condensation whereas the last one comes from matter condensation. Here

$$\frac{\partial h(T)}{\partial T} = -\frac{h(T)}{4\pi} G_2(T), \quad (16)$$

where  $G_2$  is the second Eisenstein function given by

$$G_2(T) = \frac{\pi^2}{3} - 8\pi^2 \sum_n \sigma_1(n) e^{-2\pi n T}, \quad (17)$$

and we used

$$\frac{\partial \eta(T)}{\partial T} = -\frac{\eta(T)}{4\pi} G_2(T). \quad (18)$$

On the other hand, the contribution of the matter condensates are given by

$$\frac{\partial(\log[\det A]^{1/N})}{\partial T} = -\frac{t}{4\pi N} G_2(T). \quad (19)$$

Contrary to the  $F_S$  case, this may or may not be larger than the gaugino condensate part depending on the VEV of the modulus,  $\langle T \rangle$  and the parameters  $t, N$  as we will see below in more detail.

Finally, the hidden matter condensates, through the term  $\det A$ , induce an  $F$  term in the observable matter direction,  $\phi_i$

$$F_{\phi_i} = \frac{e^{-\phi_i \phi_i^\dagger/2}}{(S + S^\dagger)^{1/2}(T + T^\dagger)^{3/2}} [\Omega(S) h(T) [\det A]^{1/N}]$$

$$\times \left( \frac{s_i}{N\phi_i} + \frac{\phi_i^\dagger}{M_v^2} \right) + (W_{3\phi_i} + K_{\phi_i} W_3)]. \quad (20)$$

This is exactly the result obtained in Ref. (5) in which the effect of matter condensation on  $F_{\phi_i}$  due to hidden matter mass terms was examined. The last two terms simply give the contribution coming from the cubic superpotential which vanishes for the solution to the F and D constraints before the hidden gauge group condensed. Since generically the F and D flat solutions give  $\langle \phi_i \rangle \sim M_v/10$  we see that for realistic values of  $s$  and  $N$  the first terms in both paranthesis in Eq. (20) (which correspond to the  $W_k$  pieces in  $F_k$ ) dominate the second ones.  $F_{\phi_i}$  obviously arises solely from matter condensation since its origin is the hidden matter mass term  $tr A\Pi$  in Eq. (4).

The  $F$  terms obtained above should be evaluated in the vacuum i.e. at the minimum of the scalar superpotential which is given by Eq. (11) (but now with  $k = S, T, \phi_i$ ) in order to find if SUSY is broken or not and in what direction in field space. This requires a complete numerical investigation of the scalar potential,  $V$  which we defer to the future since our aim is only to raise the possibility of SUSY breaking by hidden matter condensates and in the observable matter direction. Instead, we will try to answer the following general questions in the following.

(a) Can  $\langle F_S \rangle$ ,  $\langle F_T \rangle$  or  $\langle F_{\phi_i} \rangle$  be non-zero for realistic values of  $\langle S \rangle$ ,  $\langle T \rangle$  and  $\langle \phi_i \rangle$ ?

(b) Can the matter condensate contribution dominate that of the gaugino condensate in  $\langle F_S \rangle$  and/or  $\langle F_T \rangle$ ? (We remind that  $\langle F_{\phi_i} \rangle$  arises solely from matter condensates.)

(c) For what range of VEVs (and parameters  $N, s, t, r$  etc.) does one of the  $F$  terms dominate the others, e.g.  $\langle F_T \rangle \gg \langle F_{\phi_i} \rangle$  or vice versa?

From the explicit form of  $V$  in Eq. (11) it is easy to show that

$$\frac{\partial V}{\partial S} \propto \Omega_S - \frac{\Omega}{S + S^\dagger} - \Omega(\log[\det A]^{1/N})_S, \quad (21)$$

which means that  $\langle F_S \rangle = 0$  by using Eq. (13). This is the analog of the well-

known result in the pure gaugino condensate case (without the last term due to matter condensates). On the other hand, as in the pure gaugino condensate case, we find that  $\langle F_T \rangle \neq 0$  in general. With respect to  $\langle F_{\phi_i} \rangle$ , it was shown in Ref. (5) that this is always non-zero once  $W_n$  or hidden matter mass terms are taken into account. The reason is that, the  $n > 3$  terms give corrections to  $W_3$  which turn the modified F constraints into an inconsistent set of equations. Thus, the new set of F constraints cannot be solved simultaneously for any set of SM singlet scalar VEVs. Therefore, under our general assumptions, we find that  $\langle F_{\phi_i} \rangle \neq 0$  always, i.e. SUSY is always broken (by some amount which depends on the parameters of the model) in the observable matter direction in addition to the moduli direction.

We have seen that  $F_S$  and  $F_T$  have contributions from gaugino and matter condensates. What are the relative magnitudes of these two contributions? For  $F_S$  this is not a relevant question since  $\langle F_S \rangle = 0$  as we saw above. For  $F_T$  the situation is complicated since the result depends strongly on  $\langle T \rangle$  due to the term  $G_2(T)$  in Eq. (15). For large  $\langle T \rangle \sim 1$  (in units of  $M_v$ ),  $G_2(T) \sim 3$  and then the gaugino condensate part is dominant through the second term in the curly brackets in Eq. (15). Depending on  $t$  and  $N$ , the matter condensate part given by the third term in Eq. (15) may also be important. For the second and third terms to be comparable, one needs  $t \sim 12N$  which is rather large. For example in the realistic model examined in Ref. (5) with  $t = 22$  and  $N = 5$ , and gaugino condensate dominates for large  $\langle T \rangle$ . But  $G_2(T)$  is a very rapidly decreasing (increasing in absolute value which is relevant for us) function of  $T$ . A numerical analysis of  $G_2(T)$  shows that already for  $\langle T \rangle \sim 0.1$ ,  $G_2$  is large enough (in absolute value) so that the matter condensate part may become larger than the gaugino condensate part depending on  $t$  and  $N$ . In the example of Ref. (5) with  $t = 22$  and  $N = 5$ , the matter condensate part is in fact larger than the sum of the gaugino condensate contributions for  $\langle T \rangle < 0.1$ . In the intermediate range  $0.1 < \langle T \rangle < 1$  both contributions to  $\langle F_T \rangle$  are of comparable magnitude.

Finally, we would like to know when both  $\langle F_T \rangle$  and  $\langle F_{\phi_i} \rangle$  are non-zero which one dominates? This will give the direction of SUSY breaking in field space. From

Eqs. (13) and (15) we find the ratio

$$\frac{\langle F_T \rangle}{\langle F_{\phi_i} \rangle} \sim \frac{3N}{s_i} \frac{\langle \phi_i \rangle}{\langle T + T^\dagger \rangle}, \quad (22)$$

for large  $\langle T \rangle \sim 1$  and

$$\frac{\langle F_T \rangle}{\langle F_{\phi_i} \rangle} \sim \frac{t}{4\pi s_i} \langle \phi_i \rangle G_2(\langle T \rangle), \quad (23)$$

for small  $\langle T \rangle < 0.1$ . We find that (for  $\langle \phi_i \rangle \sim M_v/10$ ) in the first case  $\langle F_{\phi_i} \rangle > \langle F_T \rangle$  for  $3s_i > N$  and vice versa. For example, for  $N = 5$  if  $s_i = 1$  then  $\langle F_{\phi_i} \rangle < \langle F_T \rangle$  whereas if  $s_i = 5$  then  $\langle F_{\phi_i} \rangle > \langle F_T \rangle$ . In the second case  $\langle F_{\phi_i} \rangle > \langle F_T \rangle$  for  $s_i > 2t$  and vice versa. Note that in this case both  $\langle F_T \rangle$  and  $\langle F_{\phi_i} \rangle$  arise mainly due to matter condensates. In the example of Ref. (5),  $t = 22$  and  $s_i = 1, 5$  and therefore  $\langle F_{\phi_i} \rangle \ll \langle F_T \rangle$ .

The SUSY breaking scale in the observable sector which is given by the soft SUSY breaking masses or the gaugino mass  $m_{3/2}$ , must be phenomenologically acceptable, i.e.  $\sim O(TeV)$ . Using

$$m_{3/2} = e^{\langle K/2 \rangle} \frac{\langle W_{tot} \rangle}{M_v^2}, \quad (24)$$

and Eqs. (1), (6) and (8) for  $K$ ,  $Y^3$  and  $W_{np}$  respectively (and  $W_3$ ), this phenomenological constraint can be translated to constraints on the parameters  $N, M, r, s_i, t$  etc. These parameters of the string model not only should result in SUSY breaking either by hidden gaugino or matter condensation but also a SUSY breaking scale of  $O(TeV)$  in the observable sector. The example considered in Ref. (5) with  $N = 5$ ,  $M = 3$ ,  $r = 7$ ,  $t = 22$  and  $s_i = 1, 5$  gives  $m_{3/2} \sim 1 TeV$  as was shown by an explicit calculation.

#### 4. Conclusions and discussion

In this letter, we have shown that under quite general assumptions SUSY breaking by hidden matter condensation in the observable matter direction is possible. This should be compared with the conventional mechanism of breaking SUSY by hidden gaugino condensates and in the moduli direction. We have shown that both mechanisms are possible for a given string model. Whether one or the other occurs depends on the details of the string model such as the hidden gauge group, hidden matter content and the hidden matter mass terms and can only be decided by a detailed analysis of a given model.

In addition to the quite general assumptions we made in Section 2, a necessary condition for SUSY breaking by hidden matter condensation, with  $\langle F_{\phi_i} \rangle \neq 0$  is the following: the observable fields  $\phi_i$  whose VEVs give masses to the hidden matter must not be heavier than the hidden gauge group condensation scale,  $\Lambda_H$ . Otherwise,  $\phi_i$  decouple at  $\Lambda_H$  and SUSY can only be broken by hidden gaugino condensation. This condition puts severe constraints on the F and D flat solution at  $M_P$  since some VEVs  $\langle \phi_j \rangle$  must vanish so as not to give masses of  $O(M_v/10)$  to  $\phi_i$  from  $W_3$ . We find that for large  $\langle T \rangle \sim 1$ ,  $\langle F_{\phi_i} \rangle$  (due to matter condensates) can be either larger or smaller than  $\langle F_T \rangle$  (due to gaugino condensates) depending on  $N$  and  $s_i$ . For small  $\langle T \rangle < 0.1$  the hidden matter condensation mechanism is dominant and  $\langle F_{\phi_i} \rangle \ll \langle F_T \rangle$ . In the intermediate range  $0.1 < \langle T \rangle < 1$ ,  $\langle F_{\phi_i} \rangle \sim \langle F_T \rangle$  and matter and gaugino condensates contributions to  $\langle F_T \rangle$  are comparable.  $\tilde{F}_{\phi_i}$  arises solely from hidden matter condensation. We also find that  $\langle F_S \rangle = 0$  as in the pure gaugino condensate case.

Obviously,  $\langle T \rangle$  and the other VEVs such as  $\langle S \rangle$  and  $\langle \phi_i \rangle$  (i.e. the vacuum) are not arbitrary but are fixed dynamically by the non-perturbative superpotential  $W_{np}$ . One should minimize the scalar potential,  $V$ , given by Eq. (11) to find these VEVs in a given model. Since our aim in this work was just to show the possibility of a new SUSY breaking scenario, we did not investigate the scalar potential in detail. We have not touched upon the of the dilaton stability and cosmological constant problems which are closely connected to SUSY breaking either since this too requires a dynamical determination of the VEVs. In the future this should be

done in the framework of realistic string models such as the one investigated in Ref. (5). The vacuum which is fixed dynamically together with the parameters of the string model will determine which SUSY breaking mechanism actually occurs in a given model.

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